Realistic supersymmetric model with composite quarks

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We describe a realistic, renormalizable, supersymmetric "quindecuplet" model in which the top quark, left-handed bottom quark, and up-type Higgs boson are composite, with a compositeness scale $\sim 1-3$ TeV. The top-quark–Higgs-boson Yukawa coupling is a dynamically generated strong interaction effect, and is naturally much larger than any other Yukawa coupling. The light-quark doublets and right-handed up-type quarks are also composite but at higher energies; the hierarchy of quark masses and mixings is due to a hierarchy in the compositeness scales. Flavor-changing neutral currents are naturally suppressed, as is baryon-number violation by Planck-scale dimension-five operators. The model predicts that the most easily observable effects would be on b-quark physics and on the ρ parameter. In particular, a small negative $\Delta\rho = -\epsilon$ leads to $\Delta R_b > +2\epsilon$. There are effects on B-meson mixing and on flavor-changing neutral-current b-quark decays to leptons which might be detectable, but not on $b \rightarrow s\gamma$. The model also suggests the supersymmetry-breaking mass for the right-handed top squark might be considerably larger than that of the left-handed top squark. [S0556-2821(97)02919-6]

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I. INTRODUCTION

One of the most intriguing clues to physics beyond the standard model is the hierarchy of quark and lepton masses. In understanding why most fermions are so much lighter than the top quark and the $W^{\pm}$ and Z^0 gauge bosons, and why they seem to have a definite though ragged generational structure, we might hope to learn the mechanism of electroweak supersymmetry breaking, explain why only three generations exist, learn where Yukawa couplings come from, and obtain hints about grand-unified theory (GUT) and Planck-scale physics.

Supersymmetry provides an attractive solution to the gauge hierarchy problem, but so far has not given us an explanation for the hierarchy of quark and lepton Yukawa couplings. In many supersymmetric models the large top-quark Yukawa coupling provides the dynamics behind electroweak symmetry breaking [1], but no explanation is given for why this Yukawa coupling is so much larger than the others.

A proposal along these lines was made in [2,3] in which a dynamical mechanism for generating the top-quark mass was suggested. In this "quindecuplet" scenario, the top-quark, left-handed bottom quark, and up-type Higgs boson are part of a 15-dimensional multiplet of composite particles, each containing two "preons." The ordinary $SU(3)_c \times SU(2)_w \times U(1)_Y$ gauge interactions can be embedded into SU(5), under which the composite particles transform as $\mathbf{5}+\mathbf{10}$. The top-quark Yukawa coupling is generated by a strong-coupling effect of confinement [4], and the bottom-quark mass is generated through an effective higher-dimension operator. Viable three-generation models, employing all or part of this mechanism with the compositeness scale near to the Planck scale, were proposed in [3]. However, these models are very difficult to rule out as they have no new consequences at low energy. Furthermore, the compositeness scale cannot be scaled down to low energy as proton decay will become far too rapid.

In this paper we discuss a significantly modified version of the supersymmetric quindecuplet scenario in which the scale of compositeness of the left-handed top and bottom quarks, the right-handed top quark, and the up-type Higgs superfield can be only slightly above the weak scale, and the proton is stable. The other left-handed quark doublets and right-handed up and charm quarks are similarly composite, but are made of different preons, and are much more tightly bound. The right-handed down-type quarks, the leptons, and the down-type Higgs boson are elementary particles. A hierarchy of quark masses and mixings with a reasonable structure can be generated. Our model provides a realization of 't Hooft's idea that the Higgs boson should be composite at a scale below a few TeV and that some of the observed fermions should be composites which, due to chiral symmetry, are relatively light compared with their inverse size, with the Yukawa couplings generated via compositeness effects [5]. We have taken advantage of the recent discovery that the low-energy limit of many strongly coupled supersymmetric gauge theories contains massless composite bound states [4] (as has been anticipated for some time [6,7]).

There is vast literature on composite models of quarks and leptons, with and without supersymmetry [7]. However, we believe this example is unique in having the following features.

(1) The dynamics of the strongly coupled gauge theory we consider is tightly constrained by consistency with supersymmetry.

(2) The theory is renormalizable and weakly coupled at high energy.

(3) Many features of the hierarchy of quark masses and mixing angles may be qualitatively understood in terms of

three different compositeness scales.

(4) At the weak scale, the model is a phenomenologically viable and interesting extension of the standard model, with new strong gauge interactions at 1–3 TeV. Baryon and lepton numbers are sufficiently conserved and new sources of flavor-changing neutral currents (FCNC's) can be kept within experimental bounds.

These features make the model an ideal laboratory to study the observable effects which could arise from compositeness.

The low-energy phenomenology of the model is similar to that of the minimal supersymmetric standard model (MSSM) [8], but it has an approximate SU(6) global symmetry, of which the standard-model gauge group is a subgroup that has several interesting consequences. First, certain important low-energy signals of compositeness, including corrections to top-quark and left-handed bottom-quark couplings and to the ρ parameter, are related by SU(6) and supersymmetry. It is amusing to note that this scenario, in which the up-type Higgs and left-handed bottom quark are composite through the same dynamics, can potentially explain the reported excess in $Z\rightarrow b\bar{b}$ events[1] [9] and push the ρ parameter slightly negative without leading to other phenomenological problems, as we show in Sec. IV.

Below the confinement scale the SU(6) symmetry requires two massive supermultiplets which are not part of the MSSM—a charge 1/3 color triplet D with baryon number $-2/3$ and a charge 1 color singlet E with the quantum numbers of a proton—which we will refer to as a ''diquark'' and a ''triquark,'' respectively. These particles have ordinary gauge couplings and very small couplings to the first two generations of quarks but couple strongly to the third generation. Their masses are proportional to free parameters of the model.

To suppress flavor-changing neutral currents (FCNC's) we rely on a gauge-mediated supersymmetry-breaking scenario [11]. A viable possibility is to append another sector to the model which breaks supersymmetry and contains ''messenger'' quarks and leptons at ~ 30 TeV, as in [12], though perhaps a more compelling solution can be found. Compositeness effects change the predictions for squark masses; SU(6) relations imply that the right-handed top squark, which contains two preons carrying SU(3) color, gets a larger soft mass than the left-handed top squark, which contains one colored and one colorless preon.

Somewhat above the confinement scale it becomes possible to produce the resonances which are expected in theories with new strong interactions, which will occur as supermultiplets transforming in SU(6) representations. In analogy with QCD we guess that these will include vector bosons (plus their spin 0 and 1/2 superpartners), with quantum numbers allowing them to mix with all the ordinary SU(3)×SU(2)×U(1) gauge bosons. These resonances will enhance production of the third-family quarks and Higgs bosons at very high-energy colliders. We discuss possible experimental signals for compositeness in Sec. IV.

Some unpleasant features of the model are that we have to give up grand unification of the ordinary gauge couplings, and that the leptons have to be put in as a separate sector in order to ensure a long life for the proton.[2] Furthermore, we are required to make one dynamical assumption regarding the effects of confinement on the supersymmetry-breaking mass terms. Despite these aesthetic drawbacks, we feel this model is interesting enough to deserve study, as its features are quite different from most previous ones. In particular, some compositeness models must have a much higher compositeness scale in order to avoid problems with proton decay. Many have difficulties generating the observed hierarchy of fermion masses and mixings without also generating significant flavor-changing neutral currents, make several percent corrections to precision electroweak predictions, and/or require dynamical assumptions which are not known to be correct in any limit. This model seems to avoid all of these problems.

In the following section we describe a one-generation version of the model, and then present the full three-generation model by studying a sequence of effective field theories.

II. A MODEL OF COMPOSITE QUARKS

Our model is built around the simplest example of an $N=1$ supersymmetric gauge theory which is known to confine and to not dynamically break its global symmetries, i.e., SU(2) with chiral superfields in six doublets. This theory has an $\mathrm{SU}(6)\times\mathrm{U}(1)_R$ global chiral symmetry. By looking for a low-energy effective description of this theory which has the same global anomalies [6], moduli space of vacua and gauge-invariant operators as the high-energy theory, Seiberg was able to determine [4] that the correct low-energy effective description of this theory contains a massless gauge-singlet chiral superfield $\mathbf{M}_{ij}$ ($i,j=1,\ldots,6$,) transforming as a ''quindecuplet''—a 15-component antisymmetric tensor of the global SU(6)—interacting via the effective superpotential

$$W=Pf(\mathbf{M})=\frac{1}{6!}\,\epsilon^{ijklmn}\mathbf{M}_{ij}\mathbf{M}_{kl}\mathbf{M}_{mn}\,. \tag{2.1}$$

In our model, the dynamics behind preon confinement into quarks will be three such SU(2) supersymmetric gauge theories. The preons carry ordinary $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_w\times\mathrm{U}(1)_y$ interactions, which are embedded in the usual way into an SU(5) subgroup of the SU(6) global symmetries; a **6** branches to $(\mathbf{3},\mathbf{1},-\mathbf{1/3})+(\mathbf{1},\mathbf{2},\mathbf{1/2})+(\mathbf{1},\mathbf{1},\mathbf{O})$. The composite fields **M** include the quark doublet and up-type antiquark, an up-type Higgs doublet, a diquark, a triquark, and their superpartners. The effective superpotential (2.1) will be re-

[1]There is also a reported (less significant) deficit in $Z\rightarrow c\bar{c}$ events [9], which we cannot account for. A recent analysis suggests that a revision in charmed-meson branching fractions could account for the charm deficit in Z decay, and perhaps also affect the extraction of the $Z\rightarrow b\bar{b}$ rate [10].

[2]If all the compositeness scales of this model are taken higher than $\sim 10^{15}$ GeV as in [2,3], then we can maintain ordinary quark-lepton and gauge coupling unification. It is even possible to unify the new strong interactions with the standard gauge interactions [13]. Unfortunately, with such high confinement scales we would not find any explicit signals for compositeness.

TABLE I. Fields and symmetries.

Superfield	$SU(2)_1$	$SU(2)_2$	$SU(2)_3$	$SU(3)_c$	$SU(2)_w$	$U(1)_y$	$U(1)_B$	Z_2
$\mathbf{d}_1$	2	1	1	3	1	$-1/3$	$-1/6$	+
$\mathbf{h}_1$	2	1	1	1	2	1/2	1/2	+
$\mathbf{n}_1,\mathbf{N}_1,\mathbf{N}'_1$	2	1	1	1	1	0	$-1/2$	$-$
$\overline{\mathbf{N}}_1,\overline{\mathbf{N}}'_1$	2	1	1	1	1	0	1/2	+
$\mathbf{d}_2$	1	2	1	3	1	$-1/3$	$-1/6$	+
$\mathbf{h}_2$	1	2	1	1	2	1/2	1/2	+
$\mathbf{n}_2,\mathbf{N}_2,\mathbf{N}'_2$	1	2	1	1	1	0	$-1/2$	$-$
$\overline{\mathbf{N}}_2,\overline{\mathbf{N}}'_2$	1	2	1	1	1	0	1/2	+
$\mathbf{d}_3$	1	1	2	3	1	$-1/3$	$-1/6$	+
$\mathbf{h}_3$	1	1	2	1	2	1/2	1/2	+
$\mathbf{n}_3,\mathbf{N}_3,\mathbf{N}'_3$	1	1	2	1	1	0	$-1/2$	$-$
$\overline{\mathbf{N}}_3,\overline{\mathbf{N}}'_3$	1	1	2	1	1	0	1/2	+
$\overline{d}_i\ (i=1,2,3)$	1	1	1	$\overline{3}$	1	1/3	$-1/3$	$-$
$\overline{H}_i$	1	1	1	1	2	$-1/2$	0	+
$\overline{E}_i$	1	1	1	1	1	-1	-1	$-$
$\overline{D}_i$	1	1	1	$\overline{3}$	1	1/3	2/3	+
$\overline{e}_i$	1	1	1	1	1	1	0	$-$
l_i	1	1	1	1	2	$-1/2$	0	+

sponsible for the top-, charm, and up-quark Yukawa couplings. In order to obtain masses for the bottom, strange, and down quarks, we will need to include some additional massive particles, which are doublets under the confining SU(2) groups.

Without further ado, let us list the gauge and global quantum numbers of all chiral superfields that will appear in the model. We indicate the baryon quantum number in order to demonstrate that it is a good symmetry [aside from the $SU(2)_w$ anomaly]. Note that, as in the MSSM, B and L conservation need not be explicitly imposed, but can be accidental symmetries resulting from a combination of discrete symmetries and renormalizability. Baryon number can be guaranteed simply by imposing a discrete unbroken Z_2 R symmetry under which the superpotential changes sign. The Z_2 symmetry we choose need not guarantee lepton number conservation. Other unbroken discrete symmetries can be found which would guarantee lepton number conservation. In this paper we will simply assume lepton number is conserved for simplicity, although it would be interesting in future work to consider the consequences of allowing lepton-number violation.

As a warmup, we present a one-generation version of the model, in which the top quark gains a large Yukawa coupling and the bottom quark receives a smaller one. Consider a theory with gauge group $SU(2)\times SU(3)_c\times SU(2)_w \times U(1)_y$, where the first group factor is the confining gauge group. As matter content we take the fields in Table I with subscript 3. As a superpotential take

$$W=M\overline{\mathbf{N}}_3\mathbf{N}_3+M'\overline{\mathbf{N}}'_3\mathbf{N}'_3+\eta^E\mathbf{h}_3\mathbf{h}_3E_3+\eta^H\mathbf{h}_3\mathbf{n}_3\overline{H_3}+\eta^D\mathbf{d}_3\mathbf{n}_3\overline{D_3}+\kappa^d\mathbf{d}_3\overline{\mathbf{N}}_3\overline{d_3}+\lambda^H\mathbf{h}_3\mathbf{N}_3\overline{H_3}+\lambda^D\mathbf{d}_3\mathbf{N}_3\overline{D_3}+\lambda^e l_3\overline{e_3}\overline{H_3}. \tag{2.2}$$

Below the scale of the massive doublets the effective superpotential is

$$W=\eta^E\mathbf{h}_3\mathbf{h}_3E_3+\eta^H\mathbf{h}_3\mathbf{n}_3\overline{H_3}+\eta^D\mathbf{d}_3\mathbf{n}_3\overline{D_3}-\frac{\kappa^d\lambda^H}{M}\mathbf{d}_3\overline{d_3}\mathbf{h}_3\overline{H_3}-\frac{\kappa^d\lambda^D}{M}\mathbf{d}_3\overline{d_3}\mathbf{d}_3\overline{D_3}+\lambda^e l_3\overline{e_3}\overline{H_3}. \tag{2.3}$$

At the scale Λ the $SU(2)_3$ gauge theory becomes strong and undergoes the confinement discussed above. The six preons $\mathbf{d}_3,\mathbf{h}_3,\mathbf{n}_3$ bind into a quindecuplet containing the quark doublet $q_3\sim\mathbf{d}_3\mathbf{h}_3$, the top antiquark $\overline{u_3}\sim\mathbf{d}_3\mathbf{d}_3$, the up-type Higgs boson $H_3\sim\mathbf{h}_3\mathbf{n}_3$, and two new fields $D_3\sim\mathbf{d}_3\mathbf{n}_3$ and $E_3\sim\mathbf{h}_3\mathbf{h}_3$. The dynamical superpotential (2.1) is generated, and the resulting superpotential is

$$W=\Lambda(\eta^D\overline{E_3}E_3+\eta^H H_3\overline{H_3}+\eta^D D_3\overline{D_3})+\alpha q_3q_3D_3+\beta q_3\overline{u_3}H_3+\gamma\overline{u_3}D_3\overline{E_3}-\kappa^d\lambda^H\frac{\Lambda}{M}q_3\overline{d_3}\overline{H_3}-\kappa^d\lambda^D\frac{\Lambda}{M}\overline{u_3}\overline{d_3}\overline{D_3}+\lambda^d l_3\overline{e_3}\overline{H_3}, \tag{2.4}$$

where $\alpha\sim\beta\sim\gamma\sim1$ are introduced to account for the fact that the SU(6) symmetry which determined the superpotential (2.1) has been weakly broken by the gauge and Yukawa couplings. The fields D and E are massive; let us ignore them for the moment. The term $\beta q_3\overline{u_3}H_3$ is the top-quark Yukawa coupling; it is of order 1. The bottom-quark Yukawa coupling, $\kappa^d\lambda^H(\Lambda/M)q_3\overline{d_3}\overline{H_3}$, is naturally less than one, its exact value set by Λ/M. (The bottom quark mass also depends on the ratio $\langle H_3\rangle/\langle\overline{H_3}\rangle$.) The term $\eta^H H_3\overline{H_3}$ is the μ term (the supersymmetric mass for the Higgs bosons) which is naturally of order Λ or smaller.

Thus, for $\Lambda\sim1$ TeV, $M\sim1-40$ TeV, the model naturally generates a large top-quark mass, a smaller bottom-quark mass, and an acceptable μ term. The mass of the τ lepton is put in by hand. Two new particles D and E are massive and do not much affect physics near or below m_Z.

We now turn to the construction of the full three-generation model. The superpotential

$$W=\sum_a (M_a\overline{\mathbf{N}}_a\mathbf{N}_a+M'_a\overline{\mathbf{N}'_a}\overline{\mathbf{N}'_a})+\sum_{ai} (\eta^E_{ai}\mathbf{h}_a\mathbf{h}_aE_i + \eta^H_{ai}\mathbf{h}_a\mathbf{n}_a\overline{H_i}+\eta^D_{ai}\mathbf{d}_a\mathbf{n}_a\overline{D_i}+\kappa^d_{ai}\mathbf{d}_a\overline{\mathbf{N}}_a\overline{d_i}+\lambda^H_{ai}\mathbf{h}_a\mathbf{N}_a\overline{H_i} + \lambda^D_{ai}\mathbf{d}_a\mathbf{N}_a\overline{D_i})+\sum_{ijk}\lambda^e_{ijk}l_i\overline{e_j}\overline{H_k} \tag{2.5}$$

is the most general gauge-invariant renormalizable superpotential consistent with the global $Z_2\times U(1)_L$ symmetries, and an additional global symmetry which prevents trilinear couplings for the $\mathbf{N}'_a,\overline{\mathbf{N}'_a}$. (We forbid these latter couplings because we find they can result in unacceptable FCNC's: See Sec. II C for a variation in which lepton numbers $-1,+1$ are assigned to $\mathbf{N}'_a,\overline{\mathbf{N}'_a}$, respectively, which suppresses FCNC's, and which we hope could explain the lepton mass hierarchy.)

For convenience, we make field redefinitions so that the η coupling matrices are upper triangular and the κ matrix is lower triangular. All numerical constants except for those describing the lepton-H couplings are assumed to be of order 1.

A complete analysis of the low-energy physics of this theory follows in the next section; here we give a brief summary of the roles played by the various terms in Eq. (2.5). At each compositeness scale Λ_a, the fields E_a, H_a, and D_a combine with the composite fields $\overline{E_a}\sim\mathbf{h}_a\mathbf{h}_a$, $\overline{H_a}\sim\mathbf{h}_a\mathbf{n}_a$, $\overline{D_a}\sim\mathbf{d}_a\mathbf{n}_a$ to get masses of order $\eta_{aa}\Lambda_a$. Off-diagonal terms in the η matrices will cause these composite fields to mix slightly; the mixing angles are proportional to ratios of Λ's. The quark doubles q_a are composite fields $\mathbf{d}_a\mathbf{h}_a$, and the up-type antiquarks $\overline{u_a}$ are composite fields $\mathbf{d}_a\mathbf{d}_a$. The field H_3 will become the down-type Higgs field of the MSSM. The couplings of the down quarks to the H_3 are generated by graphs involving tree-level $\mathbf{N},\overline{\mathbf{N}}$ exchange and the matrices κ^d and Λ^H. Similar graphs, with Λ^H replaced by Λ^D, will generate couplings of the D's to up and down antiquarks. A linear combination of the composite fields $\overline{H_a}$ (which is mostly $\overline{H_3}$) will become the up-type Higgs. Its superpotential coupling to the composite quarks is generated dynamically. The couplings Λ^3_{ij3} in the last line will be responsible for lepton masses.

A. Obtaining the low-energy effective-field theory

The mode is straightforward to analyze provided that all the gauge and Yukawa couplings are weak at high energies and $M_a,M'_a\gg\Lambda_a$. (Another limit, $\Lambda_a\gg M_a$, will be briefly discussed in Sec. II C.) A realistic pattern of quark masses and mixing emerges when we assume $M_a\sim M'_a$ and take the three confining SU(2) couplings equal at short distances (with dynamical scale Λ_0). The lepton mass hierarchy is put in by hand in the superpotential. We ignore the lepton couplings for the remainder of this section.

The mass hierarchy for the quarks follows from a hierarchy among the mass terms M_a. Each $SU(2)_a$ confines at a scale $\Lambda_a\sim M_a^{2/3}\Lambda_0^{1/3}$. Our assumptions, in particular our choice of four doublets $\mathbf{N}_a,\mathbf{N}'_a,\overline{\mathbf{N}}_a,\overline{\mathbf{N}'_a}$ of approximately equal mass for each confining group, will lead to the relations for the natural order of magnitude of quark masses and mixings:

$$m_d/m_s\sim\sqrt{m_u/m_c}\sim\theta_{12}\sim(M_2/M_1)^{(1/3)},$$
$$m_s/m_b\sim\sqrt{m_c/m_t}\sim\theta_{23}\sim(M_3/M_2)^{(1/3)},$$
$$\theta_{13}\sim(M_3/M_1)^{1/3}. \tag{2.6}$$

We can choose the M_a such that these are all satisfied to within a factor of 3 experimentally.

With such a large hierarchy of scales, a step-by-step top-down effective-field theory analysis is appropriate. We ignore logarithmic effects from renormalization-group running, since these only give $O(1)$ corrections to our results.

Step I: At energy scales of order $M_1\sim M'_1(\sim3\times10^8\ \text{TeV})$, we integrate out $\mathbf{N}_1,\overline{\mathbf{N}}_1,\overline{\mathbf{N}'_1},\mathbf{N}'_1$, generating in the effective superpotential the terms

$$-\left(\frac{1}{M_1}\right)\sum_{i=1,2,3}\sum_{j=1,2,3}\mathbf{d}_1\kappa^d_{1j}\overline{d_j}(\mathbf{d}_1\lambda^D_{1i}\overline{D_i}+\mathbf{h}_1\lambda^H_{1i}\overline{H_i}). \tag{2.7}$$

The effective $SU(2)_1$ gauge theory now has six light doublets and will eventually confine.

Step II: Below the scale $M'_2\sim M_2(\sim3\times10^5\ \text{TeV})$, we integrate out $\mathbf{N}_2$, $\overline{\mathbf{N}}_2$, $\overline{\mathbf{N}'_2}$, and N'_2, inducing the superpotential terms

$$-\left(\frac{1}{M_2}\right)\sum_{i=1,2,3}\sum_{j=2,3}\mathbf{d}_2\kappa^d_{2j}\overline{d_j}(\mathbf{d}_2\lambda^D_{2i}\overline{D_i}+\mathbf{h}_2\lambda^H_{2i}\overline{H_i}). \tag{2.8}$$

Now the $SU(2)_2$ gauge theory also has six light doublets.

Step III: Below the $SU(2)_1$ confinement scale $\Lambda_1(\sim3\times10^4\ \text{TeV})$, we write down an effective theory for the composite degrees of freedom $\overline{D_1}\sim\mathbf{d}_1\mathbf{n}_1$, $\overline{E_1}\sim\mathbf{h}_1\mathbf{h}_1$, $\overline{H_1}\sim\mathbf{h}_1\mathbf{n}_1$, $q_1\sim\mathbf{d}_1\mathbf{h}_1$, and $\overline{u_1}\sim\mathbf{d}_1\mathbf{d}_1$. The dynamical couplings (2.1) are written in terms of these fields as the effective superpotential

$$\alpha_1q_1q_1\overline{D_1}+\beta_1q_1\overline{u_1}H_1+\gamma_1\overline{D_1}\,\overline{u_1}\,\overline{E_1}, \tag{2.9}$$

where α, β, and γ are of order 1 and are equal up to small SU(6)-breaking effects. For simplicity of presentation, we will set the dynamically generated α, β, γ couplings equal to 1 in Eqs. (2.10), (2.12), and (2.13). Somewhat below this scale, couplings in the original superpotential produce mass terms marrying $\overline{E_1}$ to E_1, H_1 to $\overline{H_1}$, and D_1 to $\overline{D_1}$. Only the fields q_1 (the up and down quarks) and $\overline{u_1}$ (the up antiquark) survive to low energies. We integrate out the other composite fields. The couplings induced in the effective superpotential for the light fields are

$$\sum_{i=1,2,3}\sum_{j,k=2,3}\Bigg\{-\left(\frac{\Lambda_1}{M_1}\right)\kappa^d_{1i}\overline{d_i}(q_1\lambda^H_{1j}\overline{H_j}+\overline{u_1}\lambda^D_{1j}\overline{D_j})+\left(\frac{1}{M_1}\right)\left(\frac{\lambda^H_{11}}{\eta^H_{11}}+\frac{\lambda^D_{11}}{\eta^D_{11}}\right)q_1q_1\overline{u_1}\kappa^d_{1i}\overline{d_i}-\left(\frac{1}{\Lambda_1}\right)\Bigg[\left(\frac{1}{\eta^H_{11}}\right)q_1\overline{u_1}\mathbf{h}_j(\eta^H_{j1}\mathbf{n}_j+\lambda^H_{j1}\mathbf{N}_j)$$
$$+\left(\frac{1}{\eta^D_{11}}\right)q_1q_1\mathbf{d}_j(\eta^D_{j1}\mathbf{n}_j+\lambda^D_{j1}\mathbf{N}_j)\Bigg]+\left(\frac{1}{\Lambda_1^2}\right)\left(\frac{\eta^E_{k1}}{\eta^E_{11}\eta^D_{11}}\right)\overline{u_1}\mathbf{d}_j(\eta^D_{j1}\mathbf{n}_j+\lambda^D_{j1}\mathbf{N}_j)\mathbf{h}_k\mathbf{h}_k+\left(\frac{1}{M_2\Lambda_1}\right)\mathbf{d}_2\kappa^d_{2j}\overline{d_j}\left(\frac{\lambda^H_{21}}{\eta^H_{11}}q_1\overline{u_1}\mathbf{h}_2\right.$$
$$\left.+\frac{\lambda^D_{21}}{\eta^D_{11}}q_1q_1\mathbf{d}_2\right)-\left(\frac{1}{M_2\lambda_1^2}\right)\left(\frac{\lambda^D_{21}\eta^E_{k1}}{\eta^D_{11}\eta^E_{11}}\right)\overline{u_1}\mathbf{d}_2\mathbf{d}_2\kappa^d_{2j}\overline{d_j}\mathbf{h}_k\mathbf{h}_k\Bigg\}. \tag{2.10}$$

Step IV: $\underline{\mathrm{SU}(2)}_2$ confines at $\Lambda_2\sim300$ TeV. Below this scale we rewrite the theory in terms of the light composite states $D_2,q_2,H_2,\overline{u_2},E_2$, with superpotential couplings

$$\alpha_2q_2q_2D_2+\beta_2q_2\overline{u_2}H_2+\gamma_2D_2\overline{u_2}\overline{E_2}. \tag{2.11}$$

Couplings in Eq. (2.5) result in masses for $D_2,\overline{D_2},E_2,\overline{E_2},H_2,\overline{H_2}$; integrating them out leads to superpotential terms:

$$\sum_{i=1,2,3}\sum_{j=2,3}\Bigg\{-\left(\frac{\Lambda_2}{M_2}\right)\kappa^d_{2j}\overline{d_j}(\overline{u_2}\lambda^D_{23}\overline{D_3}+q_2\lambda^H_{23}\overline{H_3})+\left(\frac{1}{M_2}\right)\left(\frac{\lambda^H_{22}}{\eta^H_{22}}+\frac{\lambda^D_{22}}{\eta^D_{22}}\right)q_2q_2\overline{u_2}\kappa^d_{2j}\overline{d_j}-\left(\frac{1}{\Lambda_2}\right)\Bigg[\left(\frac{1}{\eta^H_{22}}\right)q_2\overline{u_2}\mathbf{h}_3(\eta^H_{32}\mathbf{n}_3+\lambda^H_{32}\mathbf{N}_3)$$
$$+\left(\frac{1}{\eta^D_{22}}\right)q_2q_2\mathbf{d}_3(\eta^D_{32}\mathbf{n}_3+\lambda^D_{32}\mathbf{N}_3)\Bigg]+\left(\frac{1}{\Lambda_2^2}\right)\left(\frac{\eta^E_{32}}{\eta^E_{22}\eta^D_{22}}\right)\overline{u_2}\mathbf{d}_3(\eta^D_{32}\mathbf{n}_3+\lambda^D_{32}\mathbf{N}_3)\mathbf{h}_3\mathbf{h}_3+\left(\frac{\Lambda_1}{M_1\Lambda_2}\right)q_2\kappa^d_{1i}\overline{d_i}\left(\frac{\lambda^D_{12}}{\eta^D_{22}}q_2\overline{u_1}+\frac{\lambda^H_{12}}{\eta^H_{22}}q_1\overline{u_2}\right)$$
$$-\left(\frac{\Lambda_1}{M_1\Lambda_2^2}\right)\left(\frac{\lambda^D_{12}\eta^E_{32}}{\eta^E_{22}\eta^D_{22}}\right)\overline{u_1u_2}\kappa^d_{1i}\overline{d_i}\mathbf{h}_3\mathbf{h}_3+\left(\frac{\Lambda_2}{M_2\Lambda_1}\right)q_1\kappa^d_{2j}\overline{d_j}\left(\frac{\lambda^H_{21}}{\eta^H_{11}}q_2\overline{u_1}+\frac{\lambda^D_{21}}{\eta^D_{11}}q_1\overline{u_2}\right)\Bigg\}. \tag{2.12}$$

Step V: Below the scale $M_3'\sim M_3(\sim50\text{ TeV})$, we eliminate $\mathbf{N}_3$, $\overline{\mathbf{N}}_3$, $\overline{\mathbf{N}}_3'$, $\mathbf{N}_3'$, generating the effective superpotential terms:

$$-\left(\frac{1}{M_3}\right)\mathbf{d}_3\kappa^d_{33}\overline{d_3}(\mathbf{d}_3\lambda^D_{33}\overline{D_3}+\mathbf{h}_3\lambda^H_{33}\overline{H_3})+\left(\frac{1}{\Lambda_1M_3}\right)q_1\mathbf{d}_3\kappa^d_{33}\overline{d_3}\left(\frac{\lambda^H_{13}}{\eta^H_{11}}\overline{u_1}\mathbf{h}_3+\frac{\lambda^D_{13}}{\eta^D_{11}}q_1\mathbf{d}_3\right)-\left(\frac{1}{\Lambda_1^2M_3}\right)\left(\frac{\eta^E_{31}\lambda^D_{13}\kappa^d_{33}}{\eta^E_{11}\eta^D_{11}}\right)\overline{u_1}\overline{d_3}\mathbf{h}_3\mathbf{h}_3\mathbf{d}_3\mathbf{d}_3$$
$$-\left(\frac{1}{\Lambda_2M_3}\right)q_2\mathbf{d}_3\kappa^d_{33}\overline{d_3}\left(\frac{\lambda^H_{23}}{\eta^H_{22}}\overline{u_2}\mathbf{h}_3+\frac{\lambda^D_{23}}{\eta^D_{22}}q_2\mathbf{d}_3\right)+\left(\frac{1}{\Lambda_2^2M_3}\right)\left(\frac{\eta^E_{32}\lambda^D_{23}\kappa^d_{33}}{\eta^E_{22}\eta^D_{22}}\right)\overline{u_2}\mathbf{d}_3\mathbf{d}_3\overline{d_3}\mathbf{h}_3\mathbf{h}_3. \tag{2.13}$$

Step VI: For reasons which will be explained in Sec. III, we expect soft supersymmetry-breaking masses for scalars and gauginos, which are of order 100–1000 GeV, to be generated at a scale of about 30 TeV.

Step VII: $\mathrm{SU}(2)_3$ confines at ~1 TeV. Because the supersymmetry-breaking masses for the preons are small compared with this scale, we expect them to have little to no effect on the confining dynamics. We will make the assumption at this point that the combination of confinement and supersymmetry breaking does not give expectation values to fields carrying color. (This assumption is discussed in Sec. III C.) We write down the effective superpotential below this scale in terms of the light composite and fundamental fields:

$$W_{\text{eff}}=\Lambda_3(\eta^H_{33}H\overline{H}+\eta^D_{33}D\overline{D}+\eta^E_{33}\overline{E}E)+\alpha_3q_3q_3D+\beta_3q_3\overline{u_3}H+\gamma_3\overline{u_3}D\overline{E}-\left(\frac{\Lambda_3}{\Lambda_2}\right)\left(\frac{\eta^H_{32}\beta_2}{\eta^H_{22}}q_2\overline{u_2}H+\frac{\eta^D_{32}\alpha_2}{\eta^D_{22}}q_2q_2D\right)$$
$$-\left(\frac{\Lambda_3}{\Lambda_1}\right)\left(\frac{\eta^H_{31}\beta_1}{\eta^H_{11}}q_1\overline{u_1}H+\frac{\eta^D_{31}\alpha_1}{\eta^D_{11}}q_1q_1D\right)-\left(\frac{\Lambda_3}{M_3}\right)\kappa^d_{33}\overline{d_3}(\lambda^H_{33}q_3\overline{H}+\lambda^D_{33}\overline{u_3}\overline{D})-\left(\frac{\Lambda_2}{M_2}\right)(\kappa^d_{22}\overline{d_2}+\kappa^d_{23}\overline{d_3})(q_2\lambda^H_{23}\overline{H}+\overline{u_2}\lambda^D_{23}\overline{D})$$
$$-\left(\frac{\Lambda_1}{M_1}\right)(\kappa^d_{11}\overline{d_1}+\kappa^d_{12}\overline{d_2}+\kappa^d_{13}\overline{d_3})(q_1\lambda^H_{13}\overline{H}+\overline{u_1}\lambda^D_{13}\overline{D})+\left(\frac{\Lambda_3^2}{\Lambda_2^2}\right)\left(\frac{\eta^E_{32}\eta^D_{32}\gamma_2}{\eta^E_{22}\eta^D_{22}}\right)\overline{u_2}D\overline{E}+\left(\frac{\Lambda_3^2}{\Lambda_1^2}\right)\left(\frac{\eta^E_{31}\eta^D_{31}\gamma_1}{\eta^E_{11}\eta^D_{11}}\right)\overline{u_1}D\overline{E}$$
$$+\text{nonrenormalizable couplings}. \tag{2.14}$$

We have dropped the ''3'' subscript on the H, D, and E fields, since only one linear combination remains of each. The nonrenormalizable superpotential couplings are all suppressed by mass scales of Λ_2 or higher. A discussion of observable low-energy effects from effective nonrenormalizable terms, as well as an explanation for why we choose $\Lambda_3 \sim 1$ TeV, can be found in Sec. IV. A discussion of supersymmetry-breaking effects is in Secs. III C and III D.

Below ~ 1 TeV the model resembles the minimal supersymmetric standard model, with the addition of the massive E and D superfields. The up-quark Yukawa couplings to the up-type Higgs boson are diagonal, with the ith generation quark receiving a coupling of order $\Lambda_3/\Lambda_i \sim (M_3/M_i)^{2/3}$. The down-quark Yukawa coupling matrix is lower triangular, with the natural size of the entries in row $i \propto \Lambda_i/M_i \sim (\Lambda_0/M_i)^{1/3}$. Thus, the natural size of Cabibbo-Kobayashi-Maskawa mixing between families i and j is $\propto m_{d_i}/m_{d_j}$. This is about what is seen for the second and third families, and about a factor of 5 too small for the first and second families. There is no specific requirement on tan β, since we can adjust the overall scale of down-type Yukawa couplings by shifting the M_a. However, since the top-quark Yukawa coupling is a strong interaction effect, we do not expect that tan β will be much larger than one. The D and E couple mainly to third-family quarks; their masses, like the μ parameter $\Lambda_3 \eta_{33}^H$, are undetermined, but cannot be much above Λ_3 and certainly can be smaller. Note that all D and $\overline{D}$ couplings to quarks are of the same natural size as the quark-Higgs boson couplings and are aligned in the same basis, providing more than adequate suppression of the FCNC generated by D exchange in box diagrams.

We leave a discussion of electroweak symmetry-breaking and supersymmetry-breaking terms for Sec. III.

B. A minor variation with no strong *CP* problem

It is amusing to note that if only the second- and third-family quarks are composite, the model naturally predicts a massless up quark, which could explain the small size of strong CP violation [14]. The down-quark mass and the Cabbibo angle need not vanish. This variation can be regarded as a limiting case of the model described in the preceding section, with $\Lambda_1 \rightarrow \infty$, $M_1 \rightarrow \infty$, $\Lambda_1/M_1 \rightarrow m_d/\langle \overline{H} \rangle$.

C. When the Λ's are large

If any or all of the confining SU(2)'s become strong at a scale $\Lambda_a \gg M_a$, the effective theory analysis is very different. Seiberg has shown that the supersymmetric SU(2) gauge theory with 8 or 10 massless doublets flows to a superconformally invariant strongly interacting infrared fixed point (IRFP) [15]. We expect this to be approximately the case for our model as well when M_a, $M'_a \ll \Lambda_a$, although in the extreme infrared the masses for the doublets $\mathbf{N}_a$, $\mathbf{N}'_a$, $\overline{\mathbf{N}}_a$, $\overline{\mathbf{N}}'_a$ will push the dynamics away from the fixed point, causing the theory to confine and produce the same light particles as the limit described in the preceding section. However, in this case the theory is strongly coupled for a long momentum range above the confinement scale, whereas in the preceding section we assumed weakly coupled descriptions both above and below the confinement scale.

It is attractive to consider this regime because our superpotential has many free parameters which would be determined by properties of the IRFP. For instance, we can assign lepton number to $\mathbf{N}'_a$, $\overline{\mathbf{N}}'_a$ and add couplings of leptons to the preons, of the form $\kappa^l_{ai}\mathbf{h}_a\overline{\mathbf{N}}'_a l_i$, to Eq. (2.5). When the theory is approximately governed by the IRFP over a large energy range, the lepton and H fields acquire anomalous dimensions of order 1. Such anomalous dimensions could explain the hierarchy of lepton masses, as well as the quark masses and mixing angles, as in [16]. However, there are important subtleties involved with this idea, and it seems we cannot say anything about the theory in this limit without doing a fair bit of speculation. We leave this for a future publication.

III. BREAKING SUPERSYMMETRY AND ELECTROWEAK SYMMETRY

A. Hidden sector breaking

It is usually assumed that supersymmetry is spontaneously broken in a ''hidden'' sector, which couples only via supergravity [17]. Planck-scale physics communicates supersymmetry breaking to the visible sector, leading to apparent explicit soft supersymmetry-breaking terms. In order that squark exchange does not produce excessive flavor-changing neutral currents, it is also usually assumed that the resulting supersymmetry-breaking contribution to scalar masses is universal at the Planck scale. If squark masses are kept nearly degenerate by an approximate symmetry (which is broken only by small superpotential couplings), then a ''super Glashow-Iliopoulos-Maiani (GIM)'' mechanism prevents large FCNC's such as might contribute to the $K_L - K_S$ mass difference [18]. However Hall, Kostelecky, and Raby pointed out that the squark mass degeneracy is violated by renormalization effects below the Planck scale, and so theories which do not have approximate nonabelian flavor symmetries for the first two families may have difficulties with FCNC's [19].

A way to avoid FCNC's without squark degeneracy is to use approximate Abelian symmetries to align the squark masses with the quark masses, so that, for example, the down- and strange squark masses are diagonal in the same basis as the down- and strange quark masses [20]. Note that for the left-handed squarks, it is not possible to align both the up- and down-squark masses, since the soft supersymmetry-breaking terms are $SU(2)_w$ symmetric. Because of the small $K_L - K_S$ mass difference, it is phenomenologically necessary to align the left-handed down-squark masses rather than the left-handed up-squark masses.

In our model, there is no approximate Abelian or non-Abelian flavor symmetry for the quarks at any scale, and no reason to expect that the Planck-scale physics which communicates supersymmetry breaking should respect any such symmetry. Even if some miraculous mechanism provides degenerate squark masses at the Planck scale, the first- and second-family quarks have strong couplings of very different strengths below the Planck scale, which will induce substantial (order 1) nondegeneracy in the renormalized squark masses.

Although nondegenerate, the renormalized squark masses will tend to align with the quark masses, since the squark mass nondegeneracy is produced by the same physics responsible for the quark mass hierarchy. For the left-handed squarks, the alignment will be with the left-handed up quarks. Thus $D-\bar{D}$ mixing could be suppressed. We see no way to account for the small size of the K_L-K_S mass difference, unless the first two family squarks are very heavy ($\sim$5 TeV).

We believe the experimental absence of large FCNC's is strong evidence that if this model or any similar approach is correct, then supergravity is not the messenger of supersymmetry breaking. We must therefore look well below the Planck scale for the supersymmetry breaking and the messenger interactions.

B. New mechanisms for supersymmetry breaking in composite models

The most attractive possibility is that the same dynamics which produces the composite quarks could also result in supersymmetry breaking. Indeed many examples [21,13,22] are now known of supersymmetric theories in which gauge boson confinement, in conjunction with a superpotential, leads to dynamical supersymmetry breaking [23]. Most of these examples involve two or more gauge groups [22], and a careful analysis of the constraints following from confinement in one or more of the groups, the superpotential, and gauge D terms is required in order to uncover the dynamical supersymmetry breaking.

In the limit that all couplings except the confining SU(2)'s are turned off, our theory has a moduli space of supersymmetric ground states [4]. We have treated the superpotential terms (2.5) perturbatively, and not found any mechanism whereby these could induce supersymmetry breaking. Our model therefore appears to have the MSSM, without soft supersymmetry-breaking terms, as its low-energy limit. In our analysis so far we have neglected any dynamical effects involving $SU(3)_c$ and/or $SU(2)_w$ gauge interactions. Although it is conceivable that nonperturbative effects involving standard-model gauge groups could lead to dynamical supersymmetry breaking, the supersymmetry-breaking scale would surely be too small [24]. We therefore must modify the model in order to introduce supersymmetry breaking.

C. Gauge mediated visible sector breaking

We have outlined in the previous section why our model, like all other viable supersymmetric models, requires the addition of a ''supersymmetry-breaking sector.'' We have also explained in Sec. III A why in order to have acceptably small FCNC, supersymmetry breaking must be communicated by interactions well below $\Lambda_2\sim 300$ TeV. The possibility which is safest from FCNC is to have the ordinary SU(3)×SU(2)×U(1) interactions communicate supersymmetry breaking to the squarks and sleptons, since these interactions are flavor blind. The first two families of squarks will then naturally have sufficient degeneracy. Examples of low-energy supersymmetry-breaking sectors with gauge-mediated supersymmetry breaking have been constructed and studied elsewhere [12], and shown to be viable, with supersymmetry-breaking communicated at a scale $\sim$30 TeV. If we append such a sector to our model, the main effect will be the generation of mass terms for superpartners carrying SU(3)×SU(2)×U(1) quantum numbers, proportional to their gauge couplings squared. While it is straightforward to compute the supersymmetry-breaking masses for the scalar preons at short distances, the supersymmetry-breaking masses for the scalar $(t,b),\bar{t},D,H,E$ receive strong corrections from the strong $SU(2)_3$ dynamics. The global SU(6) symmetry can be used to predict the following approximate relations for the supersymmetry-breaking scalar mass terms $\tilde{m}^2$:

$$\tilde{m}_{\bar{t}}^2=\tilde{m}_0^2+2x\tilde{m}_d^2,\quad \tilde{m}_q^2=\tilde{m}_0^2+x(\tilde{m}_d^2+\tilde{m}_h^2),$$

$$\tilde{m}_H^2=\tilde{m}_0^2+x\tilde{m}_h^2,\quad \tilde{m}_D^2=\tilde{m}_0^2+x\tilde{m}_d^2,$$

$$\tilde{m}_{\bar{E}}^2=\tilde{m}_0^2+2x\tilde{m}_h^2, \tag{3.1}$$

where $\tilde{m}_0^2$ and x are an undetermined constants, and $\tilde{m}_d^2,\tilde{m}_h^2$ are the supersymmetry-breaking masses for the preons $\mathbf{d}_3$, $\mathbf{h}_3$ which in gauge-mediated supersymmetry-breaking scenario are expected to equal the supersymmetry-breaking masses for the d_i, l_i scalars, respectively. We expect that $x>0$ since it would be surprising for the lightest squarks to have the heaviest preons. The masses $\tilde{m}_d^2,\tilde{m}_h^2$ are the masses renormalized at an energy scale above Λ_3. SU(6) symmetry guarantees that Eq. (3.1) will survive strong renormalization effects below this scale in the long-distance effective theory, although there will be small corrections from the explicit SU(6) breaking. The large SU(6) symmetric superpotential couplings in the effective theory cause the parameters x and $\tilde{m}_0^2$ to be strongly scale dependent, with x increasing and $\tilde{m}_0^2$ decreasing at low energy.

However, it is possible that, for example, $\tilde{m}_q^2=\tilde{m}_0^2+x(\tilde{m}_d^2+\tilde{m}_h^2)$ is negative at all scales below Λ_3, and in this case color would be broken at a high scale. We make the dynamical assumption that this does not occur. If our assumption is wrong, then we must have the messenger scale of supersymmetry breaking lower than Λ_3, in which case the compositeness of the light fields will be irrelevant for supersymmetry breaking. While it may be possible to build a gauge-mediated supersymmetry-breaking model with a messenger sector near 1 TeV, it is likely that Λ_3 would even then have to be several TeV, making the model much less interesting for experiment, though no less viable.

D. Electroweak symmetry breaking

The large top-quark Yukawa coupling, in conjunction with soft supersymmetry-breaking terms, can drive electroweak symmetry breaking [1]. In our model the top Yukawa coupling is also related by the global SU(6) symmetry to the D, E couplings in Eq. (2.14), and these couplings also renormalize scalar masses. Note that Eq. (3.1) predicts that, as renormalization-group running causes $\tilde{m}_0^2$ to become negative at low energies, the first scalar mass squared to go negative is $\tilde{m}_H^2$, and so the radiative electroweak symmetry-breaking scenario is possible. In this regard the model resembles ordinary weakly coupled supersymmetry.

When the messenger scale is larger than Λ_3, it is interesting to consider the possibility that even while $\tilde{m}_q^2$ might be positive, so that color is unbroken, $\tilde{m}_H^2$ might be negative even at the confinement scale, making a radiative-breaking scenario unnecessary. In this case the model would more closely resemble technicolor or topcolor. Whether this scenario can occur (and whether color is unbroken) remains an unanswered dynamical question.

With elementary quarks and leptons, the gauge-mediated supersymmetry-breaking scenarios [12] are highly predictive and (so far) experimentally acceptable, with all supersymmetry-breaking masses determined in terms of only two parameters once the weak scale is fixed. In our model there are two undetermined strong interaction coefficients (x and $\widetilde{m}_0^2$) which affect the top- and bottom-squark masses and electroweak symmetry breaking. Thus the uncertainties due to strong $SU(2)_3$ interactions lead to reduced predictive power in this model, at least until the D and E are discovered. In particular, it is possible that the soft supersymmetry-breaking mass for H could be *larger* at the confinement scale than the mass for $\bar{H}$, due to compositeness effects. Electroweak symmetry breaking would then have to be due to a large soft supersymmetry-breaking $H-\bar{H}$ scalar mass term, and $\tan\beta=\langle H\rangle/\langle\bar{H}\rangle$ could be less than 1.

IV. EXPERIMENTAL TESTS OF QUARK COMPOSITENESS

A. Low-energy signals

First, we consider higher-dimension terms arising from the superpotential. The effective superpotential Eq. (2.14) contains higher-dimension terms involving the ordinary quarks, but these are all suppressed by high-mass scales. Since they do not give rise to FCNC at tree level, or violate any symmetries of the standard model, their effects are uninteresting at low energies. As we will argue below, the D and E fields can easily be taken too heavy or too decoupled to affect low-energy phenomena either at the tree level or through loops. All other superpotential terms are present in the minimal supersymmetric standard model and need no special analysis.

Actually, this is not quite true; there is one other set of operators we should discuss. We have prevented baryon-number violation in this model by imposing renormalizability at intermediate energies and a Z_2 symmetry. However, this does not evade the usual problem of dimension-five baryon-number-violating operators, which appear in the superpotential suppressed only by one power of M_{Planck}. As is well known [25], these operators generically lead to proton decay at far too high a rate to be consistent with experiment. Fortunately, in this *and all other low-energy fermion compositeness models* the problem is naturally solved: all such operators are suppressed by at least one factor of the confinement scale divided by M_{Planck}.

We next turn to the higher-dimension operators in the Kähler potential and those operators involving standard-model gauge fields. We search for effects of compositeness at low energies by doing an effective Lagrangian analysis. Since the confining interactions do not break supersymmetry, and since SU(6) is approximately valid at the confinement scale, we use a supersymmetric SU(6)-invariant effective Lagrangian below the compositeness scale. [Subleading effects due to soft supersymmetry-breaking and SU(6)-breaking terms could be included if desired.]

The most important corrections come from the low compositeness scale of the third-family quarks and up-type Higgs. Since q_3, $\bar{t}$, D, $\bar{E}$, and H transform as a ‘‘quindecuplet’’ chiral supermultiplet M_{ij}, the lowest dimension nonrenormalizable terms for the composite fields allowed by the global symmetries are the dimension-six operators

$$\int d^4\theta \, \frac{\mathcal{C}_1}{3\Lambda_3^2}\,[\mathrm{Tr}(\mathbf{M}^\dagger e^V \mathbf{M} e^V)]^2 + \frac{\mathcal{C}_2}{\Lambda_3^2}\left\{\mathrm{Tr}(\mathbf{M}^\dagger e^V \mathbf{M} e^V \mathbf{M}^\dagger e^V \mathbf{M} e^V)\,\frac{1}{6}\,[\mathrm{Tr}(\mathbf{M}^\dagger e^V \mathbf{M} e^V)]^2\right\}, \tag{4.1}$$

where $\mathcal{C}_{1,2}$ are unknown coefficients of order one, and e^V contains the SU(3)×SU(2)×U(1) gauge interactions necessary for standard-model gauge invariance.[3] These are of course the supersymmetric generalizations of the familiar current-current interactions.

Loop effects may also induce dimension-six terms involving ordinary SU(3)×SU(2)×U(1) gauge fields. If we use the naive dimensional analysis power counting scheme [26], which estimates the size of terms in an effective Lagrangian by using perturbation theory with the largest possible self-consistent cutoff ($4\pi\Lambda$), every additional spacetime derivative is associated with a factor $1/(4\pi\Lambda)$ and every gauge field with a factor $g/(4\pi\Lambda)$. Thus we expect q_{3L},t_R,H compositeness to induce effective operators involving the ordinary gauge fields such as

$$\int d^4\theta \, \frac{O(g^2/16\pi^2)}{\Lambda_3^2}\,\overline{D}_{\dot\alpha}(e^{-V}\overline{W^{\dot\alpha}}e^V)D_\alpha(e^V W^\alpha e^{-V}) + \text{H.c.} \tag{4.2}$$

and

$$\int d^4\theta \, \frac{O(g/16\pi^2)}{\Lambda_3^2}\,\mathbf{M}^\dagger e^V W_\alpha e^{-V} D^\alpha(e^V \mathbf{M} e^V) + \text{H.c.} \tag{4.3}$$

Because the standard model is weakly coupled at the scale Λ, these operators can be expected to be unimportant relative to Eq. (4.1).

Furthermore, since the top quark, charm quark, and up quark do not mix at all, and since the Higgs boson is not discovered and the top quark is barely studied, all effects observable now or in the near future involve the bottom quark and the expectation value of the neutral up-type Higgs boson.

Consider the $SU(4)\times SU(2)_w$ subgroup of SU(6), where $SU(3)_c$ is a subgroup of SU(4), and note that (q_3,H) transforms as a (4,2) and that the left-handed bottom quark and neutral up-type Higgs boson both have $I_3=-1/2$. It follows that the current-current interaction involving four bottom quarks, four neutral Higgs bosons, or two of each, is given

[3]We normalize $\mathbf{M}_{ij}$ through the kinetic term $\int d^4\theta \mathrm{Tr}(\mathbf{M}^\dagger e^V \mathbf{M} e^V)$. Note also that the usual definition of Λ in the compositeness literature is larger than ours by a factor of $\sqrt{4\pi}$.

by a single irreducible operator in the $I=1$ channel whose coefficient is a unique combination of $\mathcal{C}_1$ and $\mathcal{C}_2$—we have chosen our normalization of the $\mathcal{C}_i$ so that this combination is $\mathcal{C}_1+\mathcal{C}_2$. Thus, all effects involving these particles are correlated. This is a remarkable consequence of both SU(6) and supersymmetry, and it leads to interesting predictions below.

In the context of a nonsupersymmetric theory, effects of operators induced by top-quark compositeness were discussed by Georgi *et al.* [27]. They considered the effects of dimension-six operators involving the top-quark, left-handed bottom quark and gauge bosons. Their model-independent analysis found the most stringent constraint on left-handed top-quark compositeness came from the possible four-bottom-quark contribution to $B_d-\overline{B}_d$ mixing. Constraints on right-handed composite top quarks were much weaker.

Similarly, in our model, the term (4.1) includes the four-fermi interaction term

$$\left(\frac{\mathcal{C}_1+\mathcal{C}_2/4}{6\Lambda_3^2}\right)\overline{q_{3L}}\gamma^\mu q_{3L}\overline{q_{3L}}\gamma_\mu q_{3L} + \left(\frac{\mathcal{C}_2}{8\Lambda_3^2}\right)\overline{q_{3L}}\gamma^\mu\tau_a q_{3L}\overline{q_{3L}}\gamma_\mu\tau_a q_{3L}. \quad (4.4)$$

Here $q_{3L}\sim(t_L, V_{tb}b_L+V_{ts}s_L+V_{td}d_L)$. The term (4.4) gives a contribution to the $B_H^0-B_L^0$ mass difference of order

$$\Delta m_B\sim(\mathcal{C}_1+\mathcal{C}_2)\,\frac{|V_{td}^2|B_{B_d}f_B^2 m_B}{18\Lambda_3^2}, \quad (4.5)$$

which for positive $\mathcal{C}_1+\mathcal{C}_2$ has opposite sign compared to the contribution from the standard model. The value of $|V_{td}|$ extracted from B meson mixing, assuming the standard model, is close to 0.01 [28], but unitarity allows values as small as 0.004, leaving plenty of room for a large nonstandard contribution. Indeed, one can have the observed value of Δm_B with acceptable B_{B_d}, f_B, and V_{td} as long as

$$\Lambda_3>O(0.5\text{ TeV}\sqrt{\mathcal{C}_1+\mathcal{C}_2}). \quad (4.6)$$

One also needs to consider the effects of operators involving the Higgs and gauge bosons. We do not expect observable effects from operators involving the gauge field strength such as those contained in Eqs. (4.2) and (4.3), because of the $g/(16\pi^2)$ suppression factors. A strong bound on the compositeness scale comes from the operator

$$\left(\frac{\mathcal{C}_1+\mathcal{C}_2}{6\Lambda_3^2}\right)(H^\dagger i\overleftrightarrow{D}^\mu H)^2 \quad (4.7)$$

which is contained in Eq. (4.1). A general model-independent analysis of the observable effects of gauge-invariant dimension-six terms including Higgs bosons was done by Grinstein and Wise [29]. They found that the only low-energy observable resulting from the dimension-six operators with four Higgs fields and two covariant derivatives is a custodial SU(2) violating shift in the W and Z masses. Such a shift would affect the ρ parameter by an amount

$$(\Delta\rho)_{JJ}=-0.020\left(\frac{\sin\beta(\mathcal{C}_1+\mathcal{C}_2)(1\text{ TeV})^2}{\Lambda_3^2}\right), \quad (4.8)$$

where

$$\sin\beta=\frac{\langle H\rangle}{175\text{ GeV}}. \quad (4.9)$$

The constraint on $\Delta\rho$ from precision electroweak analysis [30] gives

$$\Delta\rho=-0.0015\pm0.0019^{+0.0012}_{-0.0009}, \quad (4.10)$$

where the last numbers reflect the uncertainties due to the unknown Higgs mass. In a supersymmetric model with a light Higgs boson, this should be taken to mean

$$\Delta\rho=-0.0024\pm0.0019. \quad (4.11)$$

Interesting corrections to the $Z-b-\overline{b}$ coupling come from Eq. (4.1) as well, which contains the interactions

$$\left(\frac{4\mathcal{C}_1+\mathcal{C}_2}{12\Lambda_3^2}\right)\overline{q_{3L}}\gamma_\mu q_{3L}H^\dagger i\overleftrightarrow{D}^\mu H + \left(\frac{\mathcal{C}_2}{4\Lambda_3^2}\right)\overline{q_{3L}}\gamma_\mu\tau_a q_{3L}H^\dagger i\overleftrightarrow{D}^\mu\tau_a H + \left(\frac{\mathcal{C}_1-2\mathcal{C}_2}{6\Lambda_3^2}\right)\overline{t_R}\gamma_\mu t_R H^\dagger i\overleftrightarrow{D}^\mu H. \quad (4.12)$$

The operator (4.12) can give important corrections to the top and bottom Z and W vertices, and was not considered by Georgi *et al.* The rate for $Z\to b\overline{b}$ will differ from the standard model rate. We find

$$(\gamma_b)_{JJ}\approx0.047\left(\frac{\sin^2\beta(\mathcal{C}_1+\mathcal{C}_2)(1\ \text{ TeV})^2}{\Lambda_3^2}\right), \quad (4.13)$$

where γ_b is defined by [31] $\Gamma(Z\to b\overline{b})=\Gamma^0(Z\to b\overline{b})(1+\gamma_b)$, and Γ^0 is the standard-model rate. The LEP and SLC experiments currently indicate that [9]

$$\gamma_b=0.023\pm0.007. \quad (4.14)$$

Comparison of Eqs. (4.13) and (4.8) shows that our model predicts

$$(\Delta\rho)_{JJ}=(-0.44)\sin^2\beta(\gamma_b)_{JJ}. \quad (4.15)$$

The model is potentially consistent with the results (4.14) and (4.11). If we assume the only nonstandard contributions to γ_b and $\Delta\rho$ come from compositeness, then for one-σ consistency with Eqs. (4.14) and (4.11) we must have

$$\tan\beta<1.3. \quad (4.16)$$

The left-handed top and bottom squarks give a positive contribution to $\Delta\rho$, while in our model a positive γ_b is correlated with a negative compositeness contribution to $\Delta\rho$, leading to a possible cancellation. For instance, the values $\Delta\rho=-0.0024$ and $\gamma_b=0.02$ can be consistent with $\sin\beta=1$ if the supersymmetry-breaking contribution to the scalar q_3 mass squared is $(60\text{ GeV})^2$, and left-right scalar mixing is small.

The operators (4.12) also give nonstandard flavor-changing neutral $b-s-Z$ and $b-d-Z$ couplings. Thus if the terms (4.12) account for the nonstandard γ_b measurement, the branching ratios and decay distributions for $b\rightarrow sl^+l^-$, $b\rightarrow dl^+l^-$, and $b\rightarrow s\nu\overline{\nu}$ should differ by a factor of $O(1)$ from their standard model values [28]. Experiments in the next few years will study these processes in detail.

Note that the nonstandard weak gauge boson couplings are mainly due to the composite interactions between the Higgs and $q_3,\overline{t}$. Unlike the weak gauge bosons, the photon has no Higgs component. Thus while the W and Z couplings to b and t quarks receive significant compositeness corrections, which could be as large as the standard-model one-loop corrections, the effects of compositeness on the $b\rightarrow s\gamma$ rate are smaller than the standard model contribution.

The largest nonstandard contribution to the K_L-K_S mass difference comes from the compositeness of $q_{2L}\sim(c_L,V_{cb}b_L+V_{cs}s_L+V_{cd}d_L)$, which is acceptably small provided

$$\Lambda_2>O(200\ \text{TeV}). \tag{4.17}$$

The contribution to $K-\overline{K}$ mixing from the compositeness of q_3 is smaller by a factor of $[V_{td}V_{ts}m_t/(V_{cd}V_{cs}m_c)]^2$. The compositeness-induced nonstandard couplings of q_2 and $\overline{c}$ to the weak gauge bosons are negligible.

Due to the high compositeness scale of the first family of quarks, there are no significant effects stemming from the compositeness of q_1, $\overline{u}$.

Another possible signal of top-quark compositeness comes from the top-quark Yukawa coupling. In our model the top-quark Yukawa coupling is a nonperturbative effect. A large top Yukawa coupling runs quickly towards an infrared fixed line, which typically gives in the MSSM

$$m_t\sim 200\ \text{GeV}\ \sin\beta. \tag{4.18}$$

In the MSSM $m_t/\sin\beta$ cannot be more than about 220 GeV since a large top Yukawa coupling indicates that the MSSM does not remain weakly coupled at higher energies [32]. However in our model no such bound need be satisfied.

We leave a more comprehensive analysis of the low-energy phenomenology of quark and Higgs compositeness for a future publication.

B. New particles

The model predicts that massive fields D and E must exist, in order that the composite fields of the third-generation form an SU(6) representation. Both are SU(2) singlets but are electrically charged, and D is colored. As a result they do not affect the ρ parameter and other quantities sensitive to SU(2) violation. The E does not couple to any pair of light fields, but D does, and its couplings are not flavor diagonal and can violate the GIM mechanism [33]. Fortunately for the model, the GIM mechanism applies for the first two generations, but the D has a large coupling to the third generation and can contribute to FCNC in B physics, either through loops or through direct exchange. In loop effects, limits from $b\rightarrow s\gamma$ on the scalar diquark are strongest and are similar to those on charged Higgs bosons. However, the scalar diquark receives both a supersymmetric and a supersymmetry-breaking contribution to its mass. The supersymmetry-breaking contribution could easily be larger than 500 GeV. Furthermore nothing in the model prevents us from giving the D field a large supersymmetry-preserving mass. We should therefore view $b\rightarrow s\gamma$ as a constraint which forces m_D to be large compared with m_W but moderate compared with the compositeness scale; as such it does not test the model. Exchange of D fields can induce dimension-five and -six terms which can contribute to FCNC's, but given the $b\rightarrow s\gamma$ constraint these are always subleading to the standard-model contributions. There are no significant limits on the D and E fermions or the E scalar beyond the obvious ones from collider searches. Note that although the D and E have the gauge quantum numbers of down quarks and positrons, they are forbidden to mix with them by baryon-number conservation; as a result, there will be no violations of unitarity in the CKM matrix.

Of course, nothing would substitute for the direct discovery at colliders of these supermultiplets. Generically speaking, the main decay mode of the fermionic D (the ''diquarkino''), which has odd-R parity, is to two third-family quarks and a standard-model gaugino, while the diquark decays mainly to two third-family quarks. The triquark, with even-R parity, primarily decays to three third-family quarks, and the scalar E (the ''trisquark'') decays to three third-family quarks and a gaugino. Again, the absence of mixing of D and E with down quarks and positrons distinguishes the decays of these particles from similar particles in many other models. Still, the specific decay modes depend in detail on the masses, both supersymmetry preserving and breaking, of these and other fields, and we do not have enough constraints on these masses to make definite predictions for their decay signatures.

C. Squark masses

The standard predictions of a gauge-mediated supersymmetry-breaking model with messenger quarks and leptons [12] will apply to all fundamental particles and those which are composite above the supersymmetry-breaking scale. In particular, all such quarks have roughly the same mass, with the $\text{SU}(2)_w$ doublet squarks of the first two generations being slightly heavier than the $\text{SU}(2)_w$ singlets. However, the low-energy composite fields satisfy Eq. (3.1). We expect $x\sim 1$ and $m_0^2<0$ at low energy in order that electroweak symmetry be broken. This will then lead to the prediction that $\widetilde{m}_{\overline{t}}$ is greater than $\widetilde{m}_q$ by a substantial amount, of order (very roughly) 40%. The D and E fields may have large supersymmetry-preserving masses, but were their soft supersymmetry-breaking masses to be measured, a number of simple relations, such as $\widetilde{m}^2_{\overline{t}}+\widetilde{m}^2_E\approx 2\widetilde{m}^2_q$, would be strong tests of SU(6).

D. Signals at multi-TeV colliders

The most dramatic signals of quark compositeness could be seen in collisions at energies above the scale Λ_3. Here the particle spectrum is expected to include a multitude of resonances, and the form factors for the couplings of top quarks, bottom quarks, and weak gauge bosons will differ from their standard-model values. If QCD is a good guide, the resonance region is well above the scale Λ_3, by a factor of $\sim 3-10$, and since we expect $\Lambda_3=1-3$ TeV these would probably be out of reach for LHC and any forseeable lepton collider. Still it is interesting to examine the likely high-energy signals of the new strong interactions. We expect a huge number of resonances with quite exotic quantum numbers (color sextets, weak triplets, charge two, high spins, etc), but since these will probably have a mass of $\sim 3-30$ TeV they could be out of experimental reach for the foreseeable future.

The resonances likely to have the largest effects on high-energy phenomenology are in a massive vector supermultiplet, with the quantum numbers of a 35-plet plus a singlet under the global SU(6). These have ordinary spins 1, 1/2, and 0. Their ordinary $SU(3)_c \times SU(2)_w \times U(1)$ quantum numbers are

$$(\mathbf{8},\mathbf{1},\mathbf{0})+(\mathbf{1},\mathbf{3},\mathbf{0})+\mathbf{3}(\mathbf{1},\mathbf{1},\mathbf{0})+(\mathbf{3},\mathbf{1},-\mathbf{1/3})+(\overline{\mathbf{3}},\mathbf{1},\mathbf{1/3})$$
$$+(\mathbf{1},\mathbf{2},\mathbf{1/2})+(\mathbf{1},\mathbf{2},-\mathbf{1/2})+(\mathbf{3},\mathbf{2},-\mathbf{5/6})+(\overline{\mathbf{3}},\mathbf{2},\mathbf{5/6}). \tag{4.19}$$

Their masses might well be too large to have effects at LHC, but it is worth asking how one could observe them if they are on the light side. Perhaps the largest effects at LHC could come from the heavy color-octet spin-one resonance, the analogue of the ρ meson, which mixes with the gluon and couples most strongly to the third-generation quarks. Potentially it could show up in the channels

$$gg \to t\bar{t}, \quad gg \to b\bar{b}, \quad gg \to gg. \tag{4.20}$$

Higgs boson and electroweak gauge boson production could also be enhanced through some of the other resonances and might be visible at LHC or at a lepton collider.

Clearly a more comprehensive study is needed here, which we leave for future work.

V. CONCLUSIONS

We have presented a supersymmetric quindecuplet model, using with considerable modification the mechanisms of [2,3], in which certain spin 0 and 1/2 particles of the supersymmetric standard model are composite. Our dynamical analysis relies heavily on the work of Seiberg [4]. The quark mass hierarchy is explained as a hierarchy of confinement scales, with the compositeness scale of the third generation at 1–3 TeV. The up-type Yukawa couplings are generated dynamically, the down-type Yukawas by exchange of massive fields and confinement. If gauge-mediated supersymmetry breaking is used, flavor-changing neutral currents are suppressed without fine-tuning. The model has two new supermultiplets below a TeV, and a slew of resonances well above a TeV, which couple predominantly to the third generation. An approximate global SU(6) symmetry and supersymmetry assure that confinement-scale effects on the ρ parameter, $Z \to b\bar{b}$, $b \to sl^+l^-$, $b \to dl^+l^-$, $b \to s\nu\bar{\nu}$, $b \to d\nu\bar{\nu}$ and $B-\bar{B}$ mixing are determined by $\tan\beta$ and a single unknown coefficient. The relation, Eq. (4.15), is particularly unusual and also is phenomenologically interesting given present constraints on ρ and R_b. Among the predictions which are probably generic to low-energy supersymmetric compositeness models are that the usual relations for soft supersymmetry-breaking masses are significantly modified by compositeness effects and the problem of dimension-five baryon-number-violating operators is eliminated.

While unlikely to be the full story, especially as the lepton mass hierarchy is unexplained and supersymmetry breaking requires a separate sector, this quindecuplet model has many interesting and new elements. Its ability to avoid many of the classic problems of compositeness models is remarkable. Could this be a sign that a strongly coupled supersymmetric gauge theory is indeed the missing piece of the phenomenological puzzle?

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[1] L. Ibanez and G. G. Ross, Phys. Lett. **110B**, 215 (1982); L. Alvarez-Gaume, M. Claudson, and M. B. Wise, Nucl. Phys. **B207**, 96 (1982); K. Inoue *et al.*, Prog. Theor. Phys. **68**, 927 (1982); H. P. Nilles, M. Srednicki, and D. Wyler, Phys. Lett. **120B**, 346 (1983); J. Ellis *et al.*, *ibid.* **125B**, 275 (1983); L. Alvarez-Gaume, J. Polchinski, and M. B. Wise, Nucl. Phys. **B221**, 495 (1983).

[2] A. E. Nelson (unpublished).

[3] M. J. Strassler, Phys. Lett. B **376**, 119 (1996).

[4] N. Seiberg, Phys. Rev. D **49**, 6857 (1994).

[5] G. 't Hooft, *Recent Developments In Gauge Theories*, Proceedings, Nato Advanced Study Institute, New York (Plenum, New York, 1980); see also S. Dimopoulos, S. Raby, and L. Susskind, Nucl. Phys. **B173**, 208 (1980).

[6] D. Amati, K. Konishi, Y. Meurice, and G. C. Rossi, Phys. Rep. **162**, 169 (1988).

[7] Examples of models of supersymmetric compositeness may be found in the following: R. R. Volkas and G. C. Joshi, Phys. Rep. **159**, 303 (1988), and references therein; R. D. Peccei, *The Mass Problem and the Issue of Compositeness*, Lectures given at Lake Louise Winter Institute: Selected Topics in Electroweak Interactions, Lake Louise, Canada, 1987 (Lake Louise Winter Inst., 1987), p. 564; M. Yasue, Prog. Theor. Phys. **78**, 1437 (1987); S. Yu. Khlebnikov and R. D. Peccei, Phys. Rev. D **48**, 361 (1993); J. C. Pati, *Towards a Pure Gauge Origin of the Fundamental Forces: Unity Through Preons*, Lectures given at Puri Winter School in Physics: Particle Physics and Cosmology at the Interface (BCSPIN–I.O.P.), Puri, India,

1993 (SUSY 93 Workshop), pp. 255–281; H. Terazawa, Mod. Phys. Lett. A **10**, 199 (1995); P. Dimopoulos, G. K. Leontaris, and N. D. Tracas, Report No. IOA-01-96, hep-ph/9604265.

[8] For a review see H. E. Haber and G. L. Kane, Phys. Rep. **117**, 75 (1985).

[9] The LEP Collaborations ALEPH, DELPHI, L3, OPAL, and the LEP Electroweak Working Group, CERN Report No. CERN-PRE/95-172 (unpublished); P. B. Renton, *Review of Experimental Results on Precision Tests of Electroweak Theories*, Invited talk at LP'95: International Symposium on Lepton Photon Interactions (IHEP), Beijing, People's Republic of China, 1995 (Lepton/Photon Symp., 1995), pp. 35–62; LEP Electroweak Working Group and SLD Heavy Flavor Group, Report No. LEPEWWG/96-01, 1996.

[10] I. Dunietz *et al.*, Fermilab Report No. Fermilab-pub-96/26-T, hep-ph/9606327.

[11] M. Dine and W. Fischler, Phys. Lett. **110B**, 227 (1982); C. Nappi and B. Ovrut, *ibid.* **113B**, 1751 (1982); K. Inoue *et al.*, Prog. Theor. Phys. **67**, 1889 (1982); J. Ellis, L. Ibanez, and G. G. Ross, Phys. Lett. **113B**, 283 (1982).

[12] M. Dine and A. E. Nelson, Phys. Rev. D **48**, 1277 (1993); M. Dine, A. E. Nelson, and Y. Shirman, *ibid.* **51**, 1362 (1995); M. Dine, A. E. Nelson, Y. Nir, and Y. Shirman, *ibid.* **53**, 2658 (1996).

[13] A. E. Nelson (unpublished).

[14] H. Georgi and I. N. McArthur, Harvard Report No. HUTP-81/A011, 1981 (unpublished); D. B. Kaplan and A. V. Manohar, Phys. Rev. Lett. **56**, 2004 (1986); K. Choi, C. W. Kim, and W. K. Sze, *ibid.* **61**, 794 (1988); K. Choi, Nucl. Phys. **B383**, 58 (1992); Phys. Lett. B **292**, 159 (1992); T. Banks, Y. Nir, and N. Seiberg, hep-ph/9403203.

[15] N. Seiberg, Nucl. Phys. **B435**, 129 (1995).

[16] H. Georgi, A. Nelson, and A. Manohar, Phys. Lett. **126B**, 169 (1983).

[17] See, for example, B. Ovrut and J. Wess, Phys. Lett. **112B**, 347 (1982); R. Barbieri, S. Ferrara, and D. V. Nanopoulos, *ibid.* **113B**, 219 (1982); J. Ellis and D. V. Nanopoulos, *ibid.* **116B**, 1333 (1982); H. P. Nilles, Nucl. Phys. **B217**, 366 (1983); A. H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); L. Ibanez, Phys. Lett. **118B**, 73 (1982); R. Barbieri, S. Ferrara, and C. Savoy, *ibid.* **119B**, 343 (1982); E. Cremmer, P. Fayet, and L. Giradello, *ibid.* **120B**, 41 (1983); E. Cremmer, S. Ferrara, L. Girardello, and A. Van Proeyen, Nucl. Phys. **B212**, 413 (1983).

[18] L. Hall, J. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).

[19] L. J. Hall, V. A. Kostelecky, and S. Raby, Nucl. Phys. **B267**, 415 (1986).

[20] Y. Nir and N. Seiberg, Phys. Lett. B **309**, 337 (1993).

[21] K. Intriligator, N. Seiberg, and S. H. Shenker, Phys. Lett. B **342**, 152 (1995).

[22] P. Pouliot, Phys. Lett. B **367**, 151 (1996); P. Pouliot and M. J. Strassler, *ibid.* **375**, 175 (1996); T. Kawano, Prog. Theor. Phys. **95**, 963 (1996); E. Poppitz, Y. Shadmi, and S. P. Trivedi, Nucl. Phys. **B480**, 125 (1996); Phys. Lett. B **388**, 561 (1996); C. Csaki, L. Randall, and W. Skiba, Nucl. Phys. **B479**, 65 (1996); C. Csaki, L. Randall, W. Skiba, and R. G. Leigh, Phys. Lett. B **387**, 791 (1996).

[23] E. Witten, Nucl. Phys. **B202**, 253 (1982).

[24] I. Affleck, M. Dine, and N. Seiberg, Nucl. Phys. **B256**, 557 (1985).

[25] S. Weinberg, Phys. Rev. D **26**, 287 (1982); N. Sakai and T. Yanagida, Nucl. Phys. **B197**, 533 (1982); S. Dimopoulos, S. Raby, and F. Wilczek, Phys. Lett. **112B**, 133 (1982); J. Ellis, D. V. Nanopoulos, and S. Rudaz, Nucl. Phys. **B202**, 43 (1982).

[26] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).

[27] H. Georgi, L. Kaplan, D. Morin, and A. Schenk, Phys. Rev. D **51**, 3888 (1995).

[28] For a review of the relevant *B* physics see A. Ali, DESY Report No. DESY-96-106, hep-ph/9606324.

[29] B. Grinstein and M. B. Wise, Phys. Lett. B **265**, 326 (1991).

[30] Particle Data Group, L. Montanet *et al.*, Phys. Rev. D **50**, 1173 (1994); 1995 update available on the PDG WWW pages (URL: http://pdg.lbl.gov/).

[31] G. Altarelli, R. Barbieri, and F. Caravaglios, Nucl. Phys. **B405**, 3 (1993).

[32] C. T. Hill, Phys. Rev. D **24**, 691 (1981); J. Bagger, S. Dimopoulos, and E. Masso, Phys. Rev. Lett. **55**, 1450 (1985); C. D. Frogatt, I. G. Knowles, and R. G. Moorhouse, Phys. Lett. B **249**, 273 (1990); **298**, 356 (1993); S. Dimopoulos, L. J. Hall, and S. Raby, Phys. Rev. Lett. **68**, 1984 (1992); V. Barger *et al.*, *ibid.* **68**, 3394 (1992); W. A. Bardeen *et al.*, Nucl. Phys. **B369**, 33 (1992).

[33] S. L. Glashow, J. Iliopoulous, and L. Maiani, Phys. Rev. D **2**, 1285 (1970).